Real CP Violation

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Abstract

We propose a new mechanism called "real CP violation" to originate spontaneous CP violation. Starting with a CP conserving theory with scalar fields in the adjoint representation of a global or local non-abelian symmetry, we show that even though the VEV's of such scalars are real they give rise to a spontaneous violation of CP. We provide an illustrative example of how this new mechanism of CP violation can give rise to physically significant phases which produce a complex CKM mixing matrix. This mechanism may prove useful in string models with moduli in the adjoint representation as well as in tackling the strong CP problem.

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Since its experimental discovery in 1964, many mechanisms to originate CP violation in K physics have been proposed. They can be grouped into two classes: explicit and spontaneous CP violation. In the former case the Lagrangian describing electroweak interactions contains some terms which are not CP invariant. For instance, some Yukawa couplings may be complex giving rise to a complex CKM matrix after diagonalization of the fermion mass matrices. On the contrary, in the spontaneous option one starts with a CP invariant Lagrangian, but the vacuum of the theory is not CP invariant [1]. Typically one has some scalar fields developing complex vacuum expectation values (VEV's) with some phases remaining after exploiting all the invariances of the theory. These physical phases appear in the quark mass matrices giving rise once again to a complex CKM matrix.

The possibility that CP is broken spontaneously is quite attractive. Still lacking the underlying theory explaining the origin of the Yukawa couplings, in the explicit case we introduce CP violation "by hand" in these complex couplings. Moreover if one decides that CP is not a good symmetry of the theory since the beginning, one may expect arbitrarily large violations of CP also in the strong interactions due to the presence of the θ term in the QCD Lagrangian [2]. On the other hand, the spontaneous breaking of CP by the vacuum of the theory is more linked to the "dynamics" of the theory itself and, if CP is a good symmetry to start with, the θ term has to be vanishing in the initial Lagrangian [3]. Obviously this fact does not imply by itself that the strong CP problem [4] is solved since the subsequent spontaneous violation of CP with phases in the quark mass matrices in general gives rise to an effective θ which is too large.

Here we come back to the idea that CP is broken only spontaneously. We propose a new mechanism for this breaking which does not entail the request of having complex VEV's of the scalar fields. For this reason we call it "real" CP violation and we show that it can be generally applied in theories with non-abelian global or gauge symmetries. The key-ingredient is to have a set of scalars sitting in the adjoint representation of these symmetries. Then, even if these scalars have real VEV's (which is generally the case for the real fields in the adjoint), CP is broken by the vacuum of the theory.

Apart from the interest in itself of this new mechanism for spontaneous CP violation, we think that there are potentially relevant applications. In particular it may give rise to a source of CP violation in string theories with moduli in the adjoint representation [5] and it can be relevant for the solution of the long-standing problem of strong CP violation [4]. We will elaborate more on this latter aspect in the second part ot this Letter.

First we introduce the mechanism of "real" CP violation. The way one defines CP transformations in the presence of a non-abelian symmetry presents an important difference with respect to the usual way CP is defined in the abelian case, say in QED. For simplicity, consider an SU(2) fermionic current coupled to the triplet of vector bosons W_i, $i = 1, 2, 3$. The demand that this interaction lagrangian be invariant under CP entails that W_3 and W_1 transform into themselves, while W_2 has to go into $-W_2$ under a CP transformation. This is equivalent to say that, having defined W^+ and W^- in terms of W_1 and W_2 in the usual way, CP interchanges W^+ and W^-. Consider now that we replace the W vector bosons with an SU(2) triplet of real scalar fields ϕ. Once again the presence of τ_2 in the SU(2) generators with $(\tau_2)^T = -(\tau_2)$, implies that under a CP transformation the second component ϕ_2 of the scalar triplet has to be odd if the interaction respects CP invariance. Hence, a VEV of this scalar component, although it is obviously real, leads to a spontaneous breaking of CP.

The key-point is that in the non-abelian case some of the generators are anti-symmetric and the corresponding scalar components of the adjoint representation have to be odd under CP if we want to find a consistent definition of CP to have the interaction lagrangian invariant under it.

We now come to the second task of this Letter, namely we show that the abovementioned mechanism of "real" CP violation can produce physical phases which show up at the level of the CKM quark mixing matrix. To this goal, we provide an illustrative example based on a horizontal $SU(3)_H$ symmetry which may be global or gauged. We introduce three scalar octets, that we generically denote with ϕ and a singlet ϕ_0. As for fermions, consider the 2 vector-like triplets $U_{(L,R)}$ and $D_{(L,R)}$ which are singlets under the $SU(2)$ of the standard model (SM) and triplets of the colour $SU(3)$ symmetry. They can get a direct large mass M_U and M_D, respectively. The enforcement of CP violation ensures that these masses are real. Let us now make the connection to the low-energy part of the model with the usual u and d SM quarks. Also u and d are triplets under $SU(3)_H$. Hence we can write the Yukawa couplings of the right-handed components of u and d with the left-handed components of the corresponding U and D and the above ϕ fields. Since we ask for CP conservation all these couplings are real. Notice that u_R and d_R have the same quantum numbers of U_R and D_R. Since we want to avoid that the previous Yukawa terms put into communication also the right-handed components of U and D with their left-handed counterparts, we impose a discrete symmetry under which u_R, d_R and all the ϕ fields are odd, while U and D are even. Finally we introduce also the usual SM Higgs doublet H. We now have the new couplings of H with U_R, D_R and u_L, d_L. Then, the tree level exchange of D gives rise to the effective interactions :

$$L_{eff} = \frac{\bar{d}_R(g_d\phi^a\lambda^a + g_d^{'}\phi_0)q_L H h_d}{M_D}, \tag{1}$$

where g_d, $g_d^{'}$ and h_d denote the Yukawa couplings with ϕ, ϕ_0 and H, respectively, λ^a are the Gell-Mann matrices of $SU(3)_H$, M_D is the direct mass of D and, finally, q_L is the usual doublet of the left-handed up- and down-quarks. Analogous contributions to the up quark sector arise with the different Yukawa couplings g_u, $g_u^{'}$ and h_u . When the scalar fields get a VEV, the above L_{eff} produce mass matrices for the up- and down-quarks which are hermitian. The presence of three ϕ octets assure that all components, in particular those related to the antisymmetric Gell-Mann matrices, get a nonvanishing VEV. Hence the quark mass matrices possess three phases. It is easy to see that one combination of them can never be reabsorbed by redefining the quark fields. Thus the CKM phase appears.

The fact that the quark mass matrices although complex are hermitean suggests that the "real" CP violation may prove useful in tackling the strong CP problem. Actually for the θ problem we need the full quark matrix involving both the ordinary and the heavy new quarks U and D. For instance, if we consider the down sector, we have the following renormalizable interactions and mass matrix:

$$(\bar{d}_R\bar{D}_R)\begin{pmatrix} 0 & \phi+\phi_0 \\ H & M_D \end{pmatrix}(d_L D_L)^T, \tag{2}$$

where the integration of the heavy $D_{R,L}$ fields induces L_{eff} in eq. (1). H, ϕ and ϕ_0 denote the mass terms coming from the VEV's of H, ϕ and ϕ_0, respectively. Notice that the VEV

of H can always be made real by performing a $U(1)$ hypercharge rotation on H. Given the hermiticity of the matrix block ϕ, we conclude that the determinant of the above mass matrix is real.

The θ term of the QCD lagrangian vanishes because of the initial CP invariance of the theory, while the contribution to the effective θ arising from the rotation of quark fields to bring them to the physical basis vanishes at the tree level since it is proportional to the argument of the determinant of the above quark mass matrix. The point now is that the phenomenologically required smallness of θ requires the quark mass matrix hermiticity to be spoiled by very tiny effects [3]. This computation would require the detailed formulation of a model which is beyond the scope of this Letter. Here we limit ourselves to a few comments on the possible suppression of the contributions giving rise to a nonvanishing effective θ.

The dangerous corrections spoiling the hermiticity of the quark mass matrices arise from loop contributions involving the presence of quartic terms in the ϕ fields as well as from terms of the kind $\phi^2 HH^*$. Having the scale of $SU(3)_H$ breaking large compared to the electrowek scale, the couplings of the latter terms have to be small not to create a hierarchy problem for the H mass. If we ask also for the quartic couplings to be small, we get relatively low masses for the ϕ scalars, for instance in the TeV region, where they can become accessible in next hadronic accelerators. Notice that the demand that the coefficients of the ϕ quartic terms be small may be naturally accomplished if the scale of the new physics is large enough to allow for a substantial running of such couplings.

After eliminating the danger represented by the above corrections proportional to the quartic scalar terms, we are left with only the Yukawa couplings for the fermions. However, such couplings at the loop level can induce only the renormalization of the fermion wave functions. Calling Z and Z' such wave function renormalizations of the right- and left-handed fields in eq. (2), we obtain:

$$(\bar{d_R}\bar{D_R})Z\begin{pmatrix} 0 & \phi+\phi_0 \\ H & M_D \end{pmatrix}Z'(d_L D_L)^T. \tag{3}$$

Z and Z' have to be hermitean. Hence the determinant of the whole matrix in eq. (3) remains real and, thus, there is no contribution to a non-vanishing θ from these terms.

Another possibility for suppressing the hermiticity-breaking corrections could be to supersymmetrize the proposed scheme. Then the radiative corrections to the fermion mass matrices would be suppressed by at least two powers of the ratio of the scale of low-energy SUSY breaking to the large scale of the theory. Corrections leading to fermionic wave function renormalization do not enjoy such a kind of protection, however, following the above mentioned argument, we conclude that they do not give rise to a non-vanishing θ. However, in SUSY theories the chiral supermultiplets are complex even if they belong to the adjoint representation and, hence, they may have complex VEV's in general spoiling the hermiticity of the mass matrices. We need some dynamical reason for them to take only real VEV's.

In conclusion, we have proposed the new mechanism of "real" CP violation to account for spontaneous breaking of CP in models with scalar fields in the adjoint representation of some global or local non-abelian symmetry. The mechanism allows for spontaneous CP violation even though no complex VEV occurs. The resulting CP violating phases leak to the fermionic mixing sector giving rise to a welcome complex CKM matrix. We pointed out

that this idea may find interesting applications in those string theories with moduli in the adjoint representation as well as in tackling the strong CP problem in the context of the spontaneous CP proposals. The illustrative example that we offered shows that it may be of interest to pursue in this direction to build a complete model of real CP violation in the non-SUSY or SUSY contexts.

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REFERENCES

[1] T.D. Lee, *Phys. Rev.* **D8** (1973) 1226;
R. Mohapatra and J. Pati, *Phys. Rev.* **D11** (1975) 566.

[2] G. 't Hooft, *Phys. Rev. Lett.* **37** (1976) 8;
R. Jackiw and C. Rebbi, *Phys. Rev. Lett.* **37** (1976) 172;
C. Callan, R. Dashen and D. Gross, *Phys. Lett.* **63B** (1976) 334.

[3] M. Beg and H. Tsao, *Phys. Rev. Lett.* **41** (1978) 278;
H. Georgi, *Had. J.* **1** (1978) 155;
R. Mohapatra and G. Senjanovic, *Phys. Lett.* **126B** (1978) 283;
G. Segre and H. Weldon, *Phys. Rev. Lett.* **42** (1978) 1191;
S. Barr and P. Langacker, *Phys. Rev. Lett.* **42** (1978) 1654;
A. Nelson, *Phys. Lett.* **136B** (1984) 165;
S. Barr, *Phys. Rev. Lett.* **53** (1984) 329;
S. Barr and D. Seckel, *Nucl. Phys.* **B233** (1984) 116;
S. Barr and A. Zee, *Phys. Rev. Lett.* **55** (1985) 2253;
A. Dannenberg, L. Hall and L. Randall, *Nucl. Phys.* **B271** (1986) 157;
S. Barr and A. Masiero, *Phys. Rev.* **D38** (1988) 366;
S. Barr and G. Segre, *Phys. Rev.* **D48** (1993) 302;
M. Dine, R. Leigh and A. Kagan, *Phys. Rev.* **D48** (1993) 2214.

[4] R. Peccei and H. Quinn, *Phys. Rev. Lett.***38** (1977) 1440.

[5] C. Bachas, C. Fabre and T. Yanagida, *Phys. Lett.* **B370** (1996) 49.